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x_F dependence of Drell-Yan transverse momentum broadening

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Abstract

We analyze x_F dependence of Drell-Yan transverse momentum broadening in hadron-nucleus collisions. In terms of generalized factorization theorem, we show that the x_F dependence of transverse momentum broadening, $\Delta\langle q_T^2\rangle(x_F)$, can be calculated in perturbative QCD. We demonstrate that $\Delta\langle q_T^2\rangle(x_F)$ is a good observable for studying the effects of initial-state multiple scattering and extracting quark-gluon correlation functions.

13.85.Qk, 12.38.Bx, 11.80.La, 24.85.+p

Parton multiple scattering is responsible for many interesting and important phenomena in high energy collisions involving nucleus, such as transverse momentum broadening, energy loss, as well as the nuclear suppression of quarkonium states. Understanding the mechanism of parton multiple scattering and its effects is crucial for making precise predictions of the nuclear suppression of quarkonium productions, which maybe a potential signal for the quark-gluon plasma in relativistic heavy ion collisions [1]. The parton multiple scattering can happen both at initial state or at the final state. The Drell-Yan pair production in hadron-nucleus collisions provide an excellent place to study the effects of initial state parton multiple scattering. The nuclear dependence in Drell-Yan transverse momentum spectrum for large q_T region has been studied in QCD perturbation theory [2,3]. In terms of generalized factorization theorem in QCD [4], the effects of multiple scattering can be expressed in terms of multiparton correlation functions [5], which are as fundamental as the parton distributions. However, these parton correlation functions are not well determined yet.

In this letter, we derive the x_F dependence of Drell-Yan transverse momentum broadening. We show that $\Delta\langle q_T^2\rangle(x_F)$ can be used as a good observable to study the effects of initial-state multiple scattering and parton energy loss. In addition, it can be used as an excellent observable for extracting information on multi-parton correlation functions.

Consider the Drell-Yan process in hadron-nucleus collisions, $h(p')+A(p) \to \ell^+\ell^-(q)+X$, where q is the four-momentum for the virtual photon γ^* which decays into the lepton pair. p' is the momentum for the incoming beam hadron and p is the momentum per nucleon for the nucleus with the atomic number A. Let q_T be the transverse momentum of the Drell-Yan pair, we define the averaged transverse momentum square as

$$\langle q_T^2\rangle^{hA} = \int dq_T^2 \cdot q_T^2 \cdot \frac{d\sigma_{hA}}{dQ^2 dq_T^2} \bigg/ \frac{d\sigma_{hA}}{dQ^2} \ . \tag{1}$$

In Eq. (1), Q is the total invariant mass of the lepton pair with $Q^2 = q^2$. Since single hard scattering is localized in space, only multiple scattering (at least, double scattering) are sensitive to dependence on the nuclear size (or $A^{1/3}$ type dependence). Therefore, in order to extract the effect due to multiple scattering, we introduce the nuclear enhancement of the Drell-Yan $\langle q_T^2\rangle$ as

$$\Delta\langle q_T^2\rangle \equiv \langle q_T^2\rangle^{hA} - \langle q_T^2\rangle^{hN} \ , \tag{2}$$

which is often called the transverse momentum broadening. The broadening of transverse momentum square defined in Eq. (2) should be sensitive to parton multiple scattering between nucleons inside a large nucleus.

In a perturbatively calculable hard scattering process, having an extra scattering between physically polarized partons is suppressed by a power of the hard scale [5]. Therefore, multiple scattering in momentum space between physically polarized partons correspond to an expansion in power series of $1/Q^2$. In this letter, we limit ourselves to double scattering in momentum space. With only single and double scattering, the broadening of Drell-Yan transverse momentum square, $\Delta\langle q_T^2\rangle$, can be parameterized as

$$\Delta\langle q_T^2\rangle = a + b\, A^{1/3} \ , \tag{3}$$

which is consistent to existing data [6,7]. In Eq. (3), $bA^{1/3}$-term represents the contribution directly from the double scattering which is explicitly proportional to the nuclear size ($\propto A^{1/3}$), with A the atomic weight of the nucleus target.

In Ref. [4], Qiu and Sterman argued that the factorization theorem for hadron-hadron scattering [8] should also be valid at the first non-leading power in momentum transfer, which is essential for calculating the double scattering systematically in QCD perturbation theory. According to this generalized factorization theorem, we can expand the numerator in Eq. (1) as

$$
\begin{aligned}
\int dq_T^2\, q_T^2\, \frac{d\sigma_{hA}}{dQ^2 dq_T^2} &= \sum_{a,b} \phi_{a/A}(x) \otimes C^{(0)}_{ab\to l\bar{l}}\left(x, Q^2\right) \otimes \phi_{b/h}(x') \\
&+ \frac{1}{Q^2} \sum_{a,b} \left[T_{a/A}(x) \otimes C^{(2)}\left(x, Q^2\right) \otimes \phi_{b/h}(x') \right. \\
&\qquad \left. + \phi_{a/A}(x) \otimes \bar{C}^{(2)}\left(x, Q^2\right) \otimes T_{b/h}(x') \right] + \ldots \\
&\equiv H_A^{(0)} + H_A^{(2)} + \bar{H}_A^{(2)} + \ldots \,.
\end{aligned} \tag{4}
$$

where $\otimes$ represents convolutions over partonic momenta, and "..." represents the terms that are suppressed by higher power of $1/Q^2$. In Eq. (4), $C^{(0)}$, $C^{(2)}$, and $\bar{C}^{(2)}$ are perturbatively calculable hard parts. $\phi_{b/h}(x')$ is the parton distribution of the beam hadron, and $\phi_{a/A}(x)$ is the parton distribution in the nucleus normalized by the atomic number A. $T_{a/A}(x)$ and $T_{b/h}(x')$ are the four-parton correlation functions [5] in the nucleus and the beam hadron, respectively.

Because of well-known EMC effect, as well as effect of nuclear shadowing and Fermi motion, nuclear dependence of the $\phi_{a/A}$ is nontrivial. However, if we parameterize the effective nuclear parton distribution $\phi_{a/A}$ into A^α times corresponding nucleon parton distributions, we found a very small power of α for a wide range of x [9]. Taking the EKS98 parameterization of nuclear parton distributions [10] as an example, we found that $\alpha \sim \pm 0.02 - 0.03$ for the x-range relevant to this study, which is much smaller than 1/3 for the $A^{1/3}$-type enhancements. Therefore, $H_A^{(0)}$ in Eq. (4), which is proportional to $\phi_{a/A}$ should have a very weak nuclear dependence. Similarly, $\bar{H}_A^{(2)}$ in Eq. (4) also has a weak nuclear dependence. On the other hand, the nuclear parton correlation function $T_{a/A}$ has an explicit dependence on the nuclear size ($\propto A^{1/3}$) [5,11], so as the $H_A^{(2)}$ in Eq. (4).

According to the factorization theorem [8], the denominator in Eq. (1) can also be expand in terms of power series:

$$
\begin{aligned}
\frac{d\sigma_{hA}}{dQ^2} &= \sum_{a,b} \phi_{a/A}(x, \mu^2) \otimes \frac{d\hat{\sigma}^{(0)}_{ab\to l\bar{l}}}{dQ^2}\left(x, x', \mu^2/Q^2, \alpha_s(\mu^2)\right) \otimes \phi_{b/h}(x', \mu^2) \left[1 + O\left(\frac{1}{Q^2}\right)\right] \\
&\equiv \sigma_A^{(0)} \left[1 + O\left(\frac{1}{Q^2}\right)\right] \,,
\end{aligned} \tag{5}
$$

where μ represents both renormalization and factorization scale. In Eqs. (4) and (5), all quantities are normalized by atomic number A. Substituting Eqs. (4) and (5) into Eq. (1), we obtain

$$
\langle q_T^2 \rangle^{hA} \approx \frac{H_A^{(0)} + H_A^{(2)}}{\sigma_A^{(0)}} \left[1 + O\left(\frac{A^0}{Q^2}\right)\right] \,. \tag{6}
$$

In deriving Eq. (6), we kept only terms up to $O(A^{1/3}/Q^2)$, and dropped all power correction terms without $A^{1/3}$ nuclear enhancement. Similarly, for a nucleon target, we have

$$\langle q_T^2\rangle^{hN} \approx \frac{H_N^{(0)}}{\sigma_N^{(0)}} \left[1 + O\left(\frac{A^0}{Q^2}\right)\right] . \tag{7}$$

Substituting above Eqs. (6) and (7) into our definition of the nuclear broadening of the transverse momentum square in Eq. (2), we derive

$$\Delta\langle q_T^2\rangle \approx \left[\frac{H_A^{(0)}}{\sigma_A^{(0)}} - \frac{H_N^{(0)}}{\sigma_N^{(0)}}\right] + \frac{H_A^{(2)}}{\sigma_A^{(0)}} , \tag{8}$$

where we neglected terms of $O(A^0/Q^2)$. The first term in Eq. (8) should have very weak nuclear dependence due to the fact that effective nuclear parton distributions have a very week A^α dependence. Therefore, the first term contributes to the a-term in Eq. (3). In addition, the first term should be numerically very small due to the cancellation between the two terms. On the other hand, the second term in Eq. (8) represents the double scattering contribution and it gives the $A^{1/3}$-type enhancement, and therefore, it contributes to the $bA^{1/3}$-term in Eq. (3).

At leading order, the inclusive cross section is [12]

$$\sigma_A^{(0)} \equiv \frac{d\sigma_{hA\to\ell^+\ell^-}}{dQ^2} = \sigma_0 \sum_q e_q^2 \int dx'\, \phi_{\bar{q}/h}(x') \int dx\, \phi_{q/A}(x)\, \delta(Q^2 - xx's) , \tag{9}$$

with $s = (p + p')^2$ and the Born cross section

$$\sigma_0 = \frac{4\pi\alpha_{em}^2}{9Q^2} . \tag{10}$$

Also at the leading order, Drell-Yan differential cross section is given by

$$\frac{d\sigma_{hA\to\ell^+\ell^-}}{dQ^2 dq_T^2} = \frac{d\sigma_{hA\to\ell^+\ell^-}}{dQ^2}\, \delta(q_T^2) . \tag{11}$$

Therefore, $H_A^{(0)}$ and $H_N^{(0)}$ in Eq. (8) vanish at the leading order, because $\int dq_T^2\, q_T^2\, \delta(q_T^2) = 0$, and consequently, the first term in Eq. (8), or the a-term in Eq. (3), vanishes at the leading order in α_s.

The double scattering contribution to the Drell-Yan transverse momentum broadening, $H_A^{(2)}$ in Eq. (8), was derived in Ref. [13]. At the leading order in α_s, it is given by [13]

$$H_A^{(2)} = \sigma_0 \left(\frac{4\pi^2\alpha_s}{3}\right) \sum_q \int dx'\, \phi_{\bar{q}/h}(x') \int dx\, T_{q/A}^{(I)}(x)\, \delta(Q^2 - xx's) , \tag{12}$$

where the four-parton correlation function is defined as

$$\begin{aligned} T_{q/A}^{(I)}(x) = &\int \frac{dy^-}{2\pi}\, e^{ixp^+y^-} \int \frac{dy_1^- dy_2^-}{2\pi}\, \theta(y^- - y_1^-)\, \theta(-y_2^-) \\ &\times \frac{1}{2}\, \langle p_A| F_\alpha{}^+(y_2^-) \bar{\psi}_q(0) \gamma^+ \psi_q(y^-) F^{+\alpha}(y_1^-) |p_A\rangle . \end{aligned} \tag{13}$$

These parton correlation functions are not well measured yet. By comparing the operator definition of the correlation functions and the definitions of the normal twist-2 parton distributions, Luo, Qiu, and Sterman (LQS) proposed the following model [5,11]:

$$T_{f/A}(x) = \lambda^2 A^{1/3} \phi_{f/A}(x) \ , \tag{14}$$

where λ is a free parameter to be fixed by experimental data, and was estimated in Ref. [13] as $\lambda^2 = 0.01$ GeV^2 by using the Drell-Yan data from NA10 and E772 experiment [7,14].

In order to obtain the leading order x_F-dependence of $\Delta\langle q_T^2\rangle$, we multiply $\delta(x_F - (x' - x))\, dx_F$ to the right-hand side of Eq. (12), and obtain

$$\frac{dH_A^{(0)}(x_F)}{dx_F} = \sigma_0 \left(\frac{4\pi^2\alpha_s}{3}\right) \sum_q \phi_{\bar{q}/h}(x_1)\, T_{q/A}^{(I)}(x_2)\, \frac{1}{(x_1+x_2)s} \ , \tag{15}$$

with

$$x_1 = (\sqrt{x_F^2 + 4\tau} + x_F)/2 \, , \quad \text{and} \quad x_2 = (\sqrt{x_F^2 + 4\tau} - x_F)/2 \, , \tag{16}$$

where $\tau = Q^2/s$. Combining Eqs. (8), (9), and (15), we obtain the x_F dependence of Drell-Yan transverse momentum broadening $\Delta\langle q_T^2\rangle$ at the leading order in α_s:

$$\frac{d\Delta\langle q_T^2\rangle(x_F)}{dx_F} = \left(\frac{4\pi^2\alpha_s}{3}\right) \cdot \frac{\sum_q e_q^2\, \phi_{\bar{q}/h}(x_1)\, T_{q/A}^{(I)}(x_2)/(x_1+x_2)}{\sum_q e_q^2 \int dx'\, \phi_{\bar{q}/h}(x')\, \phi_{q/A}(\tau/x)/x'} \ . \tag{17}$$

From Eq. (17), we see that $d\Delta\langle q_T^2\rangle/dx_F$ directly depends on the quark-gluon correlation functions. Therefore, the broadening is a very good observable for measuring quark-gluon correlation functions.

If we use LQS model for $T_{q/A}$ given in Eq. (14), from Eq. (17), we can derive a much simpler expression for the broadening,

$$\frac{d\Delta\langle q_T^2\rangle}{dx_F} = \left(\frac{4\pi^2\alpha_s}{3}\right) \lambda^2 A^{1/3} \ \frac{d\sigma}{dQ^2 dx_F} \Big/ \frac{d\sigma}{dQ^2} \ . \tag{18}$$

From Eq. (18), we can see that the Drell-Yan transverse momentum broadening should have similar x_F dependence as the differential cross section $d\sigma/dQ^2dx_F$ if the LQS model for $T_{q/A}^{(I)}(x)$ is valid. Therefore, by comparing the x_F dependence of $d\Delta\langle q_T^2\rangle/dx_F$ and $d\sigma/dQ^2dx_F$, we can provide an immediate test of LQS model for the correlation functions. We emphasize that even if LQS model is not a good approximation for the quark-gluon correlation functions, measuring the x_F-dependence of the nuclear broadening provides excellent information for extracting the quark-gluon correlation functions $T_{q/A}(x)$ directly, as shown in Eq. (17).

In the following, we use Eq. (18) to obtain the numerical estimates of the x_F dependence of the Drell-Yan transverse momentum broadening. Although the value of λ^2 for the correlation function is not well determined, a different value of λ^2 corresponds to a simple adjustment to the overall normalization. Therefore, the uncertainty in the value of λ^2 should not affect our following discussions.

In obtaining our following numerical results, we use the CTEQ4L distribution as the quark distributions in the nucleon. For effective quark distributions in the nucleus, we define $q_{i/A}(x) = q_{i/p}(x)R_i(x, A)$, and use EKS98 for the parameterizations of $R_i(x, A)$ [10], which fit the data well. We choose the renormalization and factorization scale to be $\mu = Q$, and choose the incoming beam energy $p = 800$ GeV which is the energy used by the Fermilab experiments [7,15–17].

In Fig. 1, we plotted $d\Delta\langle q_T^2\rangle/dx_F$ in Eq. (18) as a function of x_F with $A = 184$ and $Q = 5$ GeV and 11 GeV, respectively. The dotted lines correspond to EKS98 parameterizations of effective nuclear parton distributions. In order to separate the nuclear dependence caused by multiple scattering and that caused by the effective nuclear parton distributions, we also plotted $d\Delta\langle q_T^2\rangle/dx_F$ in solid curves with $R_i(x, A) = 1$. The difference between the solid and dotted lines is a direct consequence of the difference between nucleon and nuclear parton distributions. As shown in Eq. (17), x_2 represents the momentum fraction of a quark (or antiquark) from the nuclear target. At $x_F = 0$, we have $x_2 = Q/\sqrt{s} \approx 0.13$ for $Q = 5$ GeV, and ≈ 0.28 for $Q = 11$ GeV. It is clear that $Q = 5$ GeV and $Q = 11$ GeV cover very different range of x_2. For $Q = 5$ GeV, $x_F > 0$ corresponds to $x_2 < 0.1$ or corresponds to the shadowing region, while $x_F < 0$ covers the region of EMC effect. On the other hand, Drell-Yan pairs of $Q = 11$ GeV are not sensitive to the nuclear shadowing at all. Entire range of x_F values matches the range of EMC effect which include antishadowing region ($0.1 \leq x_2 \leq 0.2$), and EMC suppression region ($0.2 \leq x_2 \leq 0.7$), as well as Fermi motion region for larger x_2. Such dependence in effective nuclear quark distributions are clearly shown in Fig. 1. In Fig. 1a, the dotted line is above the solid line in the central region of x_F due to the fact that x_2 is in the antishadowing region. When x_F is positive, the dotted line is below the solid line because $q_{i/A}(x_2)$ is now in the shadowing region. When x_F is less than zero, the dotted line is again below the solid line due to the fact that x_2 is now in the region of EMC suppression. On the other hand, Fig. 1b shows slightly different relation between the dotted and solid line due to the fact that at $Q = 11$ GeV, x_2 covers a different range as shown in Fig. 1. Although the nuclear dependence in effective nuclear parton distributions change the x_F dependence of the Drell-Yan transverse momentum broadening, the change (e.g., the difference between the dotted and solid lines in Fig. 1) is extremely small. Therefore, the x_F dependence of the nuclear broadening, $\Delta\langle q_T^2\rangle(x_F)$ is not very sensitive to nuclear shadowing, and can be an excellent probe of parton multiple scattering.

The ratio of the cross section per nucleon from the Drell-Yan pair production in $p - A$ collisions has been used as a direct measurement of the parton energy loss in nuclear medium [15]. However, as pointed out in Ref. [15], a substantial fraction of the variation in the cross section ratios versus x_1 comes from the shadowing of $\phi_{f/A}(x_2)$ at small x_2. So that it is difficult to extract precise information in the energy loss. According to the model proposed by BDMPS [18], the transverse momentum broadening can be related to the parton energy loss [18]. One can then use the x_F dependence of the Drell-Yan transverse momentum broadening to measure the initial-state parton energy loss. As shown in Fig. 1, the variation in the shape of $\Delta\langle q_T^2\rangle(x_F)$ due to the shadowing is small. Therefore, measurement of $\Delta\langle q_T^2\rangle(x_F)$ will be more sensitive to the parton energy loss.

Cross sections on nuclear targets are often parameterized as A^α times corresponding cross sections on a nucleon target. In principle, the power α can be a function of A, q_T, x_F, as well as other physical observable. Nuclear dependence of the α is often used to study the effects of multiple scattering. In recent experimental study of the J/Ψ and Ψ' suppression in $p - A$ collisions [16], data was presented in terms of α as a function of transverse momentum

p_T in different x_F region. It was found [16] that the shapes of $\alpha(p_T)$ are very similar in different regions of x_F, and it was concluded [16] that the parton energy loss is independent of the $c\bar{c}$ energy.

Using our result of $\Delta\langle q_T^2\rangle(x_F)$, we can also estimate the α as a function of q_T in different x_F region for Drell-Yan production. We define the parameter $\alpha(q_T)$ for Drell-Yan as

$$\frac{d\sigma_{hA}}{dQ^2 dq_T^2} = A^{\alpha(q_T)} \times \frac{d\sigma_{hN}}{dQ^2 dq_T^2} \,. \tag{19}$$

If q_T is not too large, the q_T spectrum of Drell-Yan pairs can be approximately parameterized by a Gaussian form [19],

$$\frac{d\sigma_{hN}}{dQ^2 dq_T^2} \propto \frac{1}{\langle q_T^2\rangle^{hN}} \exp\left[-\frac{q_T^2}{\langle q_T^2\rangle^{hN}}\right] \,; \tag{20}$$

and

$$\frac{1}{A}\frac{d\sigma_{hA}}{dQ^2 dq_T^2} \propto \frac{1}{\langle q_T^2\rangle^{hA}} \exp\left[-\frac{q_T^2}{\langle q_T^2\rangle^{hA}}\right] \,. \tag{21}$$

Substituting Eqs. (20) and (21) into Eq. (19), we derive

$$\alpha(q_T) = 1 + \frac{1}{\ln(A)}\left[\ln\left(\frac{1}{1+\chi}\right) + \frac{\chi}{1+\chi}\frac{q_T^2}{\langle q_T^2\rangle^{hN}}\right] \,, \tag{22}$$

where $\chi \equiv \Delta\langle q_T^2\rangle / \langle q_T^2\rangle^{hN}$. In deriving Eq. (22), we used Eq. (2). In order to estimate the value of $\alpha(q_T)$, we use the value $\langle q_T^2\rangle^{hN} = 1.2$ GeV^2 for the cross section per nucleon [17]. And we use Eq. (18) to integrate over different range of x_F value to obtain the $\Delta\langle q_T^2\rangle$ in different range of x_F. For $A = 184$, we choose three different x_F ranges, which are the same as the regions used in Ref. [16], small x_F (SXF: $-0.1 \leq x_F \leq 0.3$), intermediate x_F (IXF: $0.2 \leq x_F \leq 0.6$), and large x_F (LXF: $0.3 \leq x_F \leq 0.93$). For these three x_F regions, we obtain corresponding values of $\Delta\langle q_T^2\rangle$ as [0.10, 0.043, 0.026] GeV^2 for $Q = 5$ GeV, and [0.078, 0.043, 0.027] GeV^2 for $Q = 11$ GeV, respectively. It is clear that the value of χ should be very small, and the power $\alpha(q_T)$ in Eq. (22) can be approximated as

$$\alpha(q_T) \approx 1 + \frac{\chi}{\ln(A)}\left[-1 + \frac{q_T^2}{\langle q_T^2\rangle^{hN}}\right] \,. \tag{23}$$

From Eq. (23), we conclude that if q_T is not too large, $\alpha(q_T)$ should show linear dependence on q_T^2 (or quadratic dependence on q_T). This feature should be universal for any Gaussian-like transverse momentum distributions. For Drell-Yan production, we expect an extremely small coefficient due to a weak transverse momentum broadening (or a small χ value in Eq. (23)).

In the above discussion, we used the same averaged value of $\langle q_T^2\rangle^{hN}$ for different x_F range. In principle $\langle q_T^2\rangle^{hN}$ should also have x_F dependence due to kinematics. The larger $|x_F|$, the smaller $\langle q_T^2\rangle^{hN}$. Hence $\langle q_T^2\rangle^{hN}$ should have similar x_F dependence as $\Delta\langle q_T^2\rangle(x_F)$ shown in Fig. 1. Therefore, $\chi = \Delta\langle q_T^2\rangle/\langle q_T^2\rangle^{hN}$ should have even smaller dependence on x_F.

In Fig. 2, we plotted our predictions for the value of $\alpha(q_T)$ defined in Eq. (22) as a function of q_T in three different x_F regions. From Fig. 2, it is clear that the shape of

$\alpha(q_T)$ in different x_F range are similar to what was observed in recent data on J/Ψ and Ψ' suppression [16]. Such similarity is natural because both q_T spectrum of Drell-Yan and Charmonium production can be approximated by a Gaussian form when q_T is not too large. In addition, we emphasize that as shown in Figs. 1 and 2, even though $\alpha(q_T)$ seems not to be sensitive to x_F, the x_F dependence of transverse momentum broadening, $\Delta\langle q_T^2\rangle(x_F)$, can be very strong.

In summary, we derived the x_F dependence of the Drell-Yan transverse momentum broadening in terms of quark-gluon correlation functions. We demonstrated that $\Delta\langle q_T^2\rangle(x_F)$ has a strong x_F dependence. In particular, if the LQS model for the correlation functions is valid, $\Delta\langle q_T^2\rangle(x_F)$ should have similar x_F dependence as the differential cross section $d\sigma/dQ^2dx_F$. We also demonstrated that $\Delta\langle q_T^2\rangle(x_F)$ is a very sensitive observable to study the effects of initial state multiple scattering and the parton energy loss. In fact, $\Delta\langle q_T^2\rangle(x_F)$ itself is an excellent observable for extracting direct information on the nuclear quark-gluon correlation functions.

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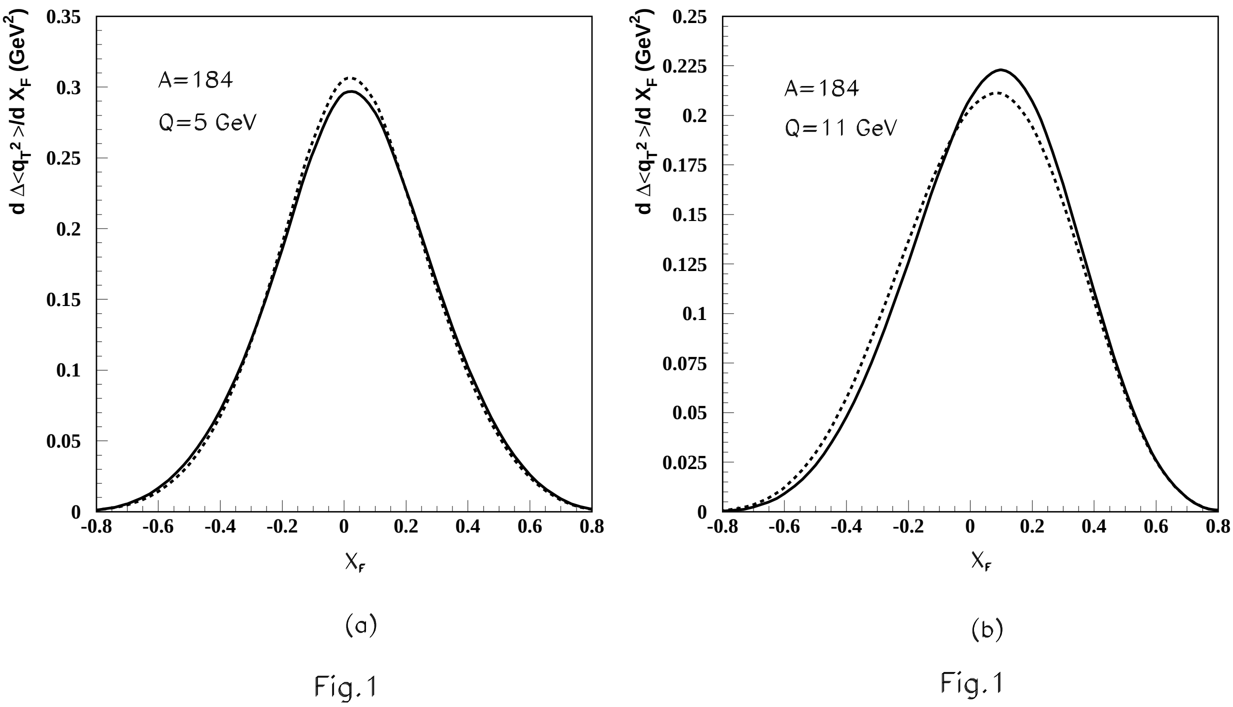


FIG. 1. The transverse momentum broadening of the Drell-Yan pair, $d\Delta\langle q_T^2\rangle/dx_F$, as a function of x_F for 800 GeV proton beam on nuclear target A=184, at $Q = 5$ GeV (a) and $Q = 11$ GeV (b).

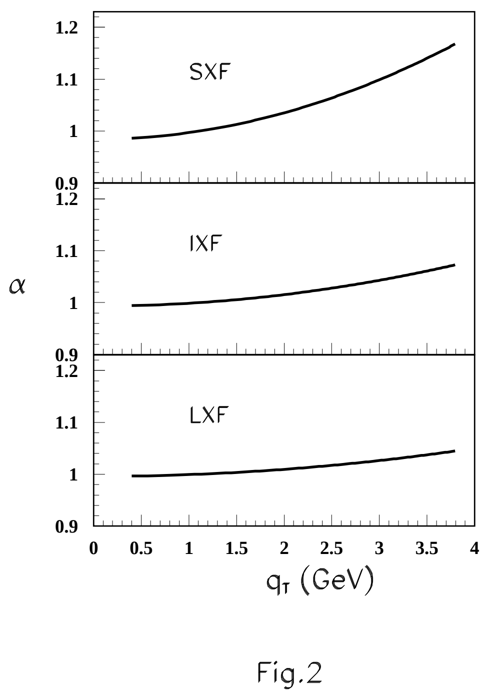
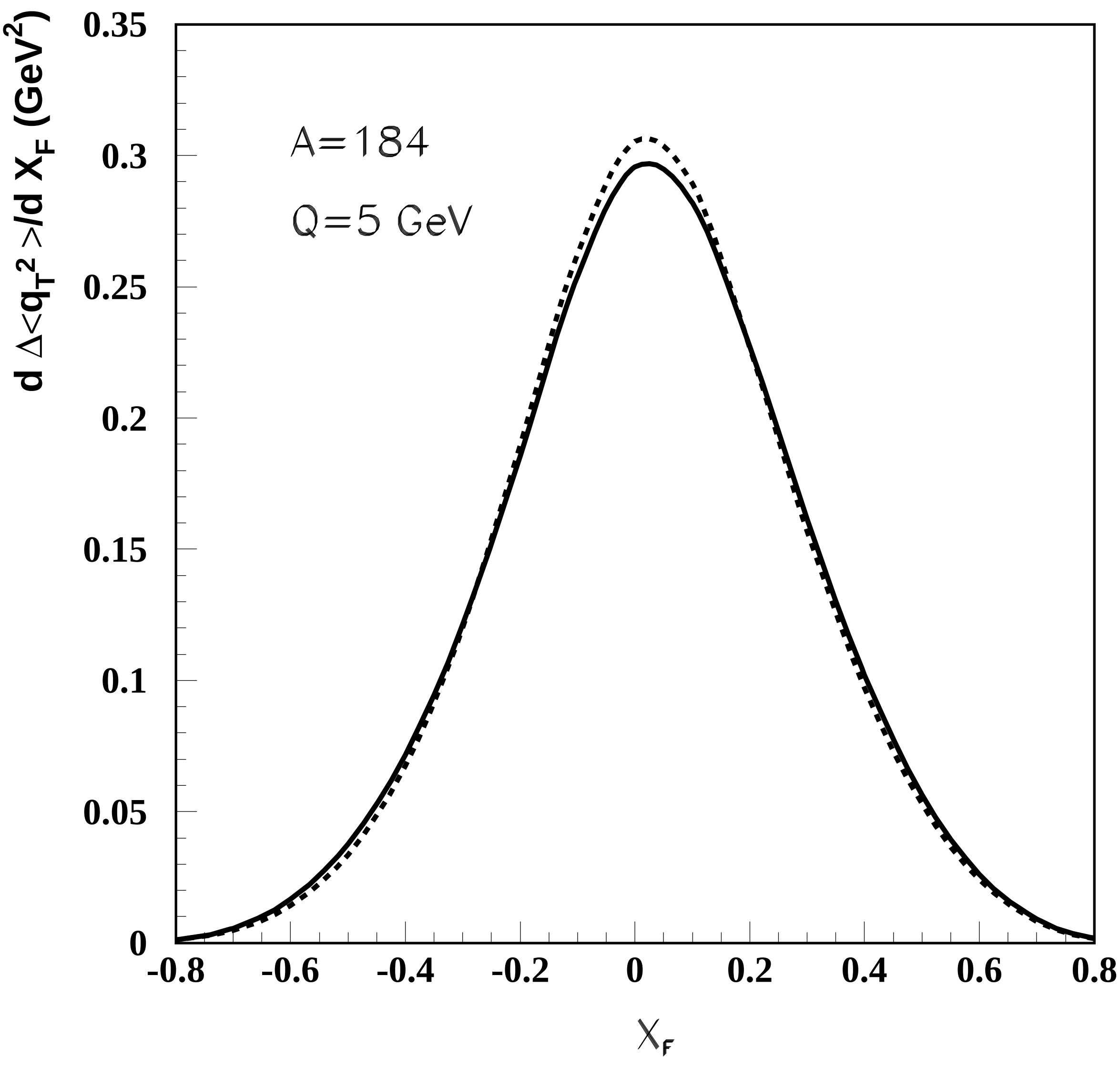


FIG. 2. The $\alpha(q_T)$ as a function of q_T for small, intermediate, and large x_F region, respectively. The curves are for $Q = 5$ GeV and 800 GeV beam energy.

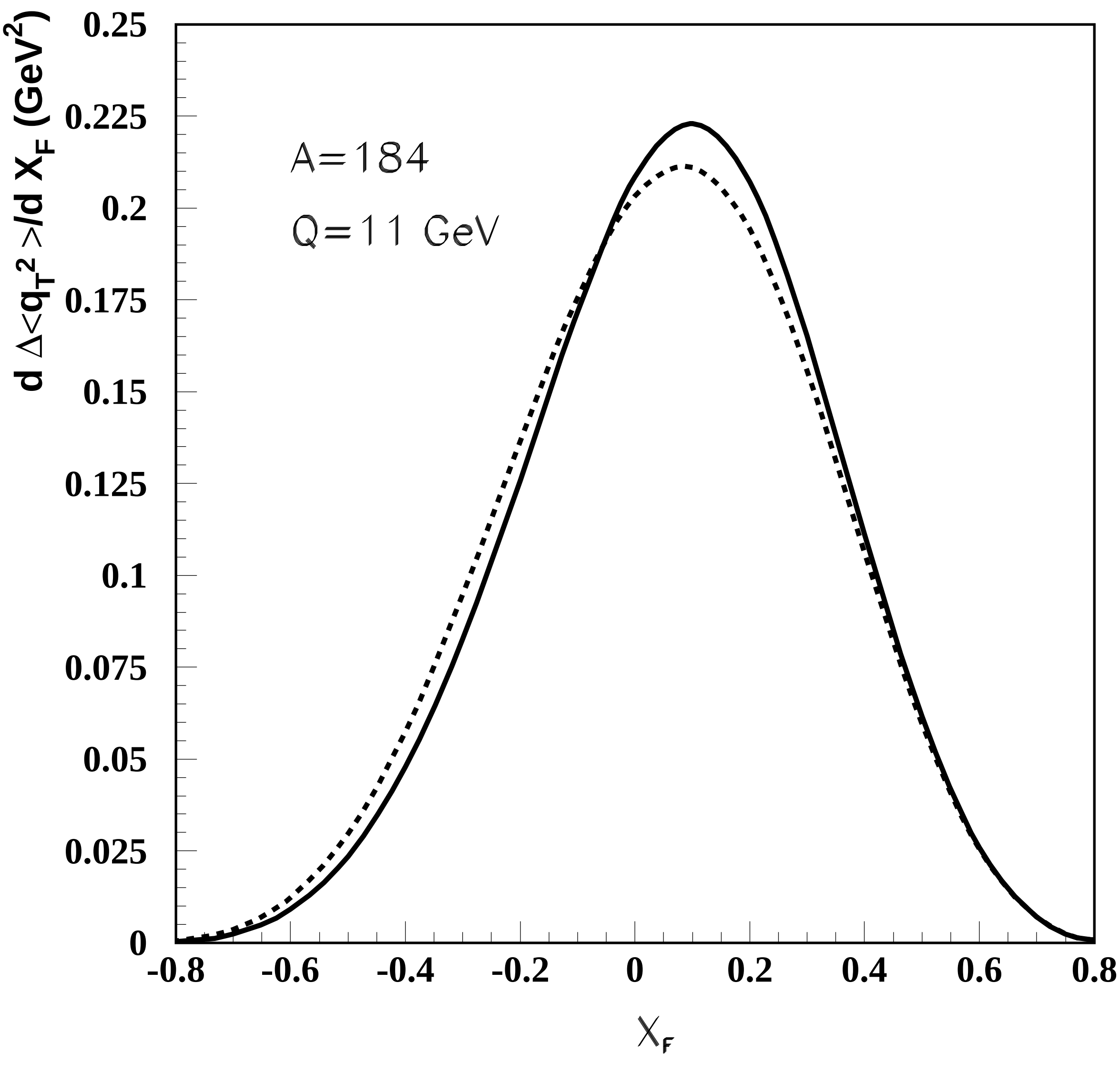
0 0.5 1 1.5 2 2.5 3 3.5 4



A=184
Q=5 GeV
d Δ<q_T^2>/d X_F (GeV^2)
0.35
0.3
0.25
0.2
0.15
0.1
0.05
0
-0.8
-0.6
-0.4
-0.2
0
0.2
0.4
0.6
0.8
X_F

(a)

Fig.1



d Δ<q_T^2>/d X_F (GeV^2)
0.25
0.225
0.2
0.175
0.15
0.125
0.1
0.075
0.05
0.025
0
A=184
Q=11 GeV
-0.8
-0.6
-0.4
-0.2
0
0.2
0.4
0.6
0.8
X_F

(b)

Fig.1

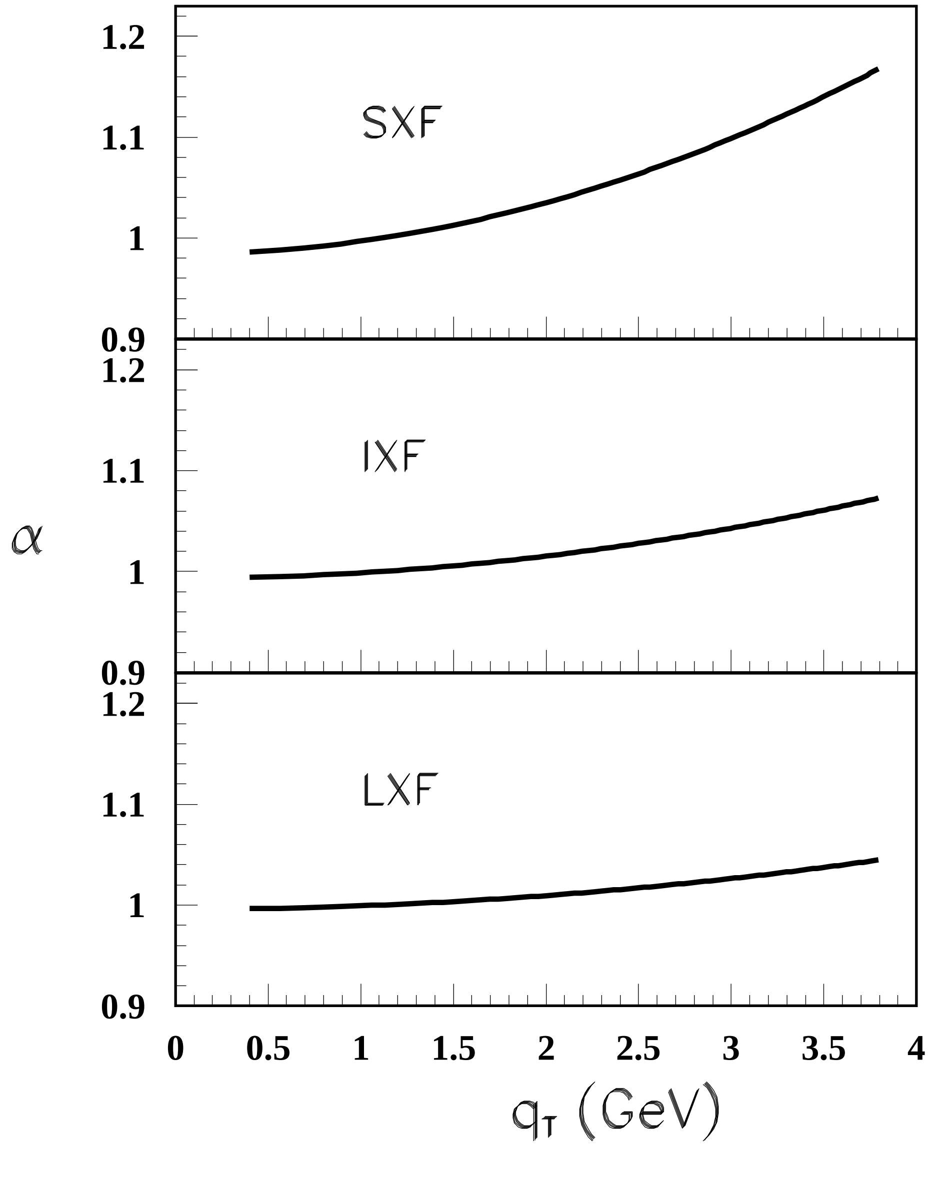


Fig.2